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Large Top Mass and Non-Linear Representation of Flavour Symmetry

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We consider an effective theory (ET) approach to flavour-violating processes beyond the Standard Model (SM), where the breaking of flavour symmetry is described by spurion fields whose low-energy vacuum expectation values are identified with the SM Yukawa couplings. Insisting on canonical mass dimensions for the spurion fields, the large top-quark Yukawa coupling also implies a large expectation value for the associated spurion, which breaks part of the flavour symmetry already at the UV scale Λ of the ET. Below that scale, flavour symmetry in the ET is represented in a non-linear way by introducing Goldstone modes for the partly broken flavour symmetry and spurion fields transforming under the residual symmetry. As a result, the dominance of certain flavour structures in rare quark decays can be understood in terms of the $1/\Lambda$ expansion in the ET. We also discuss the generalization to 2-Higgs–doublet models with large $\tan\beta$.

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Understanding the origin behind the existence of three quark and lepton families is one of the biggest challenges in contemporary particle physics. While many arguments have been given that the explanation of the electroweak symmetry breaking in the SM requires new physics at the TeV scale, the recent experiments at B-meson factories seem to indicate that quark flavour violation beyond the known sources in the Standard Model (SM) happens at much higher scales. For this reason it has been proposed to consider the SM and its potential extensions as an effective theory (ET), where the concept of quark (and also lepton) flavour is introduced by a symmetry principle that guarantees minimal flavour violation (MFV), i.e. any new source for quark flavour transitions should be induced by the Yukawa coupling matrices for up- and down-type quarks of the SM [1]. As we will explain below, in the original formulation of MFV the large top-quark Yukawa coupling distorts the conventional power-counting in the ET. In this letter, we show how the special role of the top quark can be implemented into the ET approach in a natural way, which also leads to some new insights into the structure of MFV.

As is well known [2], the maximal quark flavour symmetry group commuting with the SM gauge symmetry is

$$G_F = SU(3)_{Q_L} \times SU(3)_{U_R} \times SU(3)_{D_R} \,. \qquad (1)$$

Additional $U(1)$ factors, which become relevant for the discussion of 2-Higgs-doublet models [1], can be identified with baryon number, hyper-charge and the Peccei-Quinn symmetry [3]. The quarks are grouped into left-handed doublets and right-handed singlets with respect to weak isospin. Under the flavour symmetry G_F, they transform as

$$Q_L \sim (3,1,1)\,, \quad U_R \sim (1,3,1)\,, \quad D_R \sim (1,1,3)\,, \qquad (2)$$

while the SM gauge and Higgs fields are singlets with respect to G_F.

Considering the SM as the leading (dimension-4) part of an effective theory, the breaking of the flavour group can be achieved by introducing spurion fields [1], transforming under G_F as

$$Y_U \sim (3,\bar{3},1)\,, \qquad Y_D \sim (3,1,\bar{3})\,. \qquad (3)$$

If we assume the spurions to have canonical mass dimension one, the Yukawa couplings in the SM, which couple left- and right-handed quark fields to the Higgs doublet, already stem from dimension-five operators,

$$-\mathcal{L}_{\rm Yuk} = \frac{1}{\Lambda}(\bar{Q}_L\tilde{\phi})Y_U U_R + \frac{1}{\Lambda}(\bar{Q}_L\phi)Y_D D_R + \text{h.c.} \qquad (4)$$

where ϕ is the usual Higgs doublet, $\tilde{\phi}$ its charge conjugate, and Λ denotes some high-energy scale, $\Lambda \gg M_W$.

If the spurion fields Y_U and Y_D acquire vacuum expectation values at an intermediate scale $\Lambda' \ll \Lambda$, one could easily explain the smallness of the u, c, d, s, b quark masses, while the top quark would not fit into this scheme and requires $\Lambda' \sim \Lambda$. In other words, the top-quark Yukawa coupling should rather be described by a dimension-4 operator in the ET below Λ. In the usual (linear) formulation of MFV [1], the unusually large value of y_t thus distorts the conventional power-counting of the spurion analysis. Therefore, it would be desirable to single out the role of the top quark more explicitly. To this end, we are going to construct a non-linear representation of the flavour symmetry group in which the concept of MFV can be embedded in an economic way.

Starting from the flavour group G_F, which is considered a good symmetry at scales above Λ, we now assume that at the scale Λ the spurion Y_U acquires a vacuum expectation value (VEV) [8],

$$\langle \hat{Y}_U \rangle \equiv \langle \frac{Y_U}{\Lambda} \rangle = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & y_t \end{pmatrix}\,, \qquad (5)$$

which breaks the original flavour group G_F down to

$$G'_F = SU(2)_{Q_L} \times SU(2)_{U_R} \times SU(3)_{D_R} \times U(1)_T \,. \quad (6)$$

The reduced flavour group G'_F corresponds to the limit where all Yukawa couplings except for y_t and all off-diagonal entries in the CKM matrix are neglected. The 7 unbroken generators of the subgroup $SU(2)_{Q_L} \times SU(2)_{U_R} \times U(1)_T$ of G'_F are identified as $T^{a=1,2,3}_{Q_L,U_R}$, and a charge operator,

$$Q_T = B - \frac{2}{\sqrt{3}}\, T^8_{Q_L=U_R} \,, \quad (7)$$

which acts on the 3rd-family component of Q_L and U_R, only (here B denotes baryon number). The "Goldstone" modes, corresponding to the 9 broken generators of G_F/G'_F, are written in the standard parameterization [4],

$$\mathcal{U}(\Pi_X) = \exp\left(\frac{i}{\Lambda}\sum_{a=4}^{8} T^a_X\, \Pi^a_X\right) , \quad (8)$$

with $X = Q_L, U_R$ and the constraint $\Pi^8_{Q_L} = -\Pi^8_{U_R}$ (the case $\Pi^8_{Q_L} = \Pi^8_{U_R}$ is included in (7)).

The remaining degrees of freedom in the Yukawa matrix $\hat{Y}_U$ can be parameterized as

$$\hat{Y}_U = \mathcal{U}(\Pi_{Q_L}) \begin{pmatrix} Y_U^{(2)}/\Lambda & \begin{matrix} 0 \\ 0 \end{matrix} \\ 0 \;\; 0 & y_t \end{pmatrix} \mathcal{U}^\dagger(\Pi_{Q_R}) \,. \quad (9)$$

where the residual spurion $Y_U^{(2)}$ now has canonical mass dimension one, and transforms as $(2,2,1)_0$ under G'_F. It is supposed to develop its VEV at a lower scale, $\Lambda' \ll \Lambda$, which breaks $SU(2)_{Q_L} \times SU(2)_{Q_R}$. The 8 real parameters of the complex 2×2 matrix $Y_U^{(2)}$, the top-Yukawa coupling y_t, and the 9 Goldstone degrees of freedom add up to 18 real parameters describing the complex 3×3 matrix Y_U.

The matrices $\mathcal{U}(\Pi_{Q_L})$ and $\mathcal{U}(\Pi_{U_R})$ can be further reduced, according to

$$(\mathcal{U}(\Pi_X)_{ij} = (\Xi_X)_i\, \delta_{j3} + \sum_{k=1,2} (\mathcal{U}_X^{(2)})_{ik}\, \delta_{kj} \,, \quad (10)$$

with the individual components transforming as 3-vectors and 3×2 matrices, respectively,

$$\begin{aligned} \Xi_X(\Pi_X) &\to \Xi_X(\Pi'_X) = V_X\, \Xi_X(\Pi_X) \,, \\ \mathcal{U}_X^{(2)}(\Pi_X) &\to \mathcal{U}_X^{(2)}(\Pi'_X) = V_X\, \mathcal{U}_X^{(2)}(\Pi_X)\, V_X^{(2)\dagger} \,, \end{aligned} \quad (11)$$

with $V_X \in SU(3)_X$ and $V_X^{(2)} \in SU(2)_X$. Notice that in general, Π'_X is a complicated non-linear function of Π_X. However, on the subgroup G'_F of G_F, the Π_X transform linearly, according to

$$(\Pi'_X)^a = \mathrm{tr}\left[T^a_X\, V_X^{(2)}\, \Pi_X\, V_X^{(2)\dagger}\right] . \quad (12)$$

The unitarity of $\mathcal{U}$ implies the following relations

$$\sum_{j=1,2} (\mathcal{U}_X^{(2)})_{ij}\, (\mathcal{U}_X^{(2)\dagger})_{jk} + (\Xi_X)_i\, (\Xi_X^\dagger)_k = \delta_{ik} \,, \quad (13)$$

and

$$\begin{aligned} &\mathcal{U}_X^{(2)\dagger}\, \mathcal{U}_X^{(2)} = \mathbf{1} \,, \quad \Xi_X^\dagger\, \Xi_X = 1 \,, \\ &\mathcal{U}_X^{(2)\dagger}\, \Xi_X = \Xi_X^\dagger\, \mathcal{U}_X^{(2)} = 0 \,. \end{aligned} \quad (14)$$

The SM Yukawa term for up-type quarks in the non-linear representation can then be written in manifestly invariant form as

$$\begin{aligned} -\mathcal{L}^{(u)}_{\rm yuk} &= y_t\, (\bar{Q}_L\, \Xi_{Q_L})\, \tilde{\phi}\, (\Xi^\dagger_{U_R} U_R) \\ &\quad + \frac{1}{\Lambda}\, (\bar{Q}_L\, \mathcal{U}_{Q_L}^{(2)}\; \tilde{\phi}\; Y_U^{(2)}\, \mathcal{U}_{U_R}^{(2)\dagger}\, U_R) + \text{h.c.} \end{aligned} \quad (15)$$

Similarly, the Yukawa terms for down-type quarks are obtained from

$$\begin{aligned} -\mathcal{L}^{(d)}_{\rm yuk} &= \frac{1}{\Lambda}\, (\bar{Q}_L\, \Xi_{Q_L})\, \phi\, (\xi^\dagger_{D_R} D_R) \\ &\quad + \frac{1}{\Lambda}\, (\bar{Q}_L\, \mathcal{U}_{Q_L}^{(2)}\; \phi\; Y_D^{(2)}\, D_R) + \text{h.c.} \end{aligned} \quad (16)$$

where we have identified two irreducible spurions of G'_F in addition to $Y_U^{(2)}$ in (9),

$$\begin{aligned} \xi^\dagger_{D_R} &= \Xi^\dagger_{Q_L}\, \hat{Y}_D\, \Lambda \;\sim\; (1,1,\bar{3})_{+1} \,, \\ Y_D^{(2)} &= \mathcal{U}_{Q_L}^{(2)\dagger}\, \hat{Y}_D\, \Lambda \;\sim\; (2,1,\bar{3})_0 \,. \end{aligned} \quad (17)$$

Again, the new spurion fields have canonical mass dimension, and assume VEVs of the order $\Lambda' \ll \Lambda$. Therefore, by construction, in the limit $\Lambda \to \infty$, only the dimension-4 term for the top Yukawa coupling (first line in (15)) survives, while the dimension-5 operators in (15,16), related to the breaking of G'_F, vanish.

Noticing that the spurion matrices $Y_U^{(2)}$ and Y_D can – as usual – be diagonalized by bi-unitary transformations,

$$\begin{aligned} Y_U^{(2)} &= V_{u_L}^{(2)}\, Y_U^{(2)\rm diag}\, V_{u_R}^{(2)} \,, \\ Y_D &= V_{d_L}\, Y_D^{\rm diag}\, V_{d_R} \,, \end{aligned} \quad (18)$$

the CKM matrix in the SM is identified as

$$V_{\rm CKM} = \begin{pmatrix} V_{u_L}^{(2)\dagger}\, \mathcal{U}_{Q_L}^{(2)\dagger} \\ \Xi^\dagger_{Q_L} \end{pmatrix} V_{d_L} \,. \quad (19)$$

The construction of MFV operators in the ET is now straightforward. Let us consider the decay $b \to s\gamma$ as an example. In the SM, it is generated by flavour-changing neutral currents, arising from loop diagrams with charged gauge bosons and up-type quarks. In the ET, new physics can contribute via higher-dimensional operators above the electro-weak scale. An example is given by

$$\mathcal{O}_{\rm eff} = \frac{1}{\Lambda^3}\, (\bar{Q}_L \Xi_{Q_L}\, \phi\, \sigma_{\mu\nu}\; \xi^\dagger_{D_R} D_R)\, F^{\mu\nu} + \text{h.c.} \quad (20)$$

which is manifestly invariant under gauge and flavour transformations (a second flavour structure involving $Y_D^{(2)}$ is related to (20) and a flavour-diagonal operator by the unitarity relation (13)). After changing to the mass basis, using (18,19), we obtain

$$\left(V_{d_L}^\dagger\, \Xi_{Q_L}\, \xi_{D_R}^\dagger V_{d_R}\right)_{ij} = V_{ti}^* V_{tj}\, (y_d)_j\, \Lambda\,. \tag{21}$$

This may be compared with the standard (linear) formulation of MFV [1], where the effective operator would be written as

$$\mathcal{O}_{\rm eff} = \frac{1}{\Lambda^2}\, (\bar Q_L\, \hat Y_U \hat Y_U^\dagger\, \phi\, \sigma_{\mu\nu}\, \hat Y_D\, D_R)\, F^{\mu\nu} + \text{h.c.} \tag{22}$$

with the resulting flavour coefficients proportional to

$$\sum_{k=u,c,t} V_{ki}^*\, (y_k)^2\, V_{kj}\, (y_d)_j\,, \tag{23}$$

which coincides with (21) up to terms of order y_c^2/y_t^2. Inserting the Higgs VEV into (20) and projecting on $b \to s$ transitions, we obtain contributions to the low-energy weak effective Hamiltonian [5] with flavour coefficients

$$\begin{aligned} &\frac{m_W^2}{\Lambda^2}\, m_b\, V_{ts}^* V_{tb} \qquad (\text{for } b_R \to s_L \text{ transitions}) \\ &\frac{m_W^2}{\Lambda^2}\, m_s\, V_{tb}^* V_{ts} \qquad (\text{for } b_L \to s_R \text{ transitions}) \end{aligned} \tag{24}$$

Our example exhibits the usual advantage of the non-linear representation for the power-counting in ETs with spontaneously broken global symmetries: The dominance of the top-quark contribution in (23) is a manifest consequence of the $1/\Lambda$ expansion. Also the chiral suppression factors m_b, m_s can be traced back to the additional power of $1/\Lambda$ in (20) compared to (22), i.e. in the non-linear representation NP contributions to $b \to s\gamma$ are related to dimension-7 operators in the ET above the electro-weak scale.

Quark bi-linears with other chirality structures can be made invariant under the flavour group in a similar fashion as in the above example. The possible structures fall into three classes, which can be constructed in terms of fundamental building blocks which transform as singlets under $SU(2)_{U_R} \times SU(3)_{D_R}$. The first class contains all possible combinations of the form

$$\begin{pmatrix} \bar Q_L\, \Xi_{Q_L} \\ \bar U_R\, \Xi_{U_R} \\ \bar D_R\, \xi_{D_R} \end{pmatrix} \otimes \begin{pmatrix} \Xi_{Q_L}^\dagger\, Q_L \\ \Xi_{U_R}^\dagger U_R \\ \xi_{D_R}^\dagger D_R \end{pmatrix}. \tag{25}$$

The second class is constructed from

$$\begin{pmatrix} \bar Q_L\, \mathcal{U}_{Q_L}^{(2)} \\ \bar U_R\, \mathcal{U}_{U_R}^{(2)}\, Y_U^{(2)\dagger} \\ \bar D_R\, Y_D^{(2)\dagger} \end{pmatrix} \otimes \mathcal{P} \otimes \begin{pmatrix} \mathcal{U}_{Q_L}^{(2)\dagger} Q_L \\ Y_U^{(2)}\, \mathcal{U}_{U_R}^{(2)\dagger} U_R \\ Y_D^{(2)}\, D_R \end{pmatrix}, \tag{26}$$

where $\mathcal{P}$ is an arbitrary polynomial of $Y_U^{(2)}\, Y_U^{(2)\dagger}$ and $Y_D^{(2)}\, Y_D^{(2)\dagger}$, which transforms as $(1,1,1)_0 + (3,1,1)_0$. Finally, the third class involves

$$\begin{pmatrix} \bar Q_L\, \mathcal{U}_{Q_L}^{(2)} \\ \bar U_R\, \mathcal{U}_{U_R}^{(2)}\, Y_U^{(2)\dagger} \\ \bar D_R\, Y_D^{(2)\dagger} \end{pmatrix} \otimes \mathcal{P} \cdot Y_D^{(2)}\, \xi_{D_R} \otimes \begin{pmatrix} \Xi_{Q_L}^\dagger\, Q_L \\ \Xi_{U_R}^\dagger U_R \\ \xi_{D_R}^\dagger D_R \end{pmatrix} + \text{h.c.} \tag{27}$$

Our discussion can be generalized to models with an extended Higgs sector. Of particular interest is the 2-Higgs–doublet models (2HDM) of type II [6], where the small value of the bottom quark mass can be related to a large ratio of Higgs VEVs $\tan\beta = v_u/v_d$, such that

$$y_b = m_b/v_d \sim y_t = m_t/v_u \sim 1\,.$$

The non-linear representation of the flavour symmetry for this case can be achieved by considering a VEV for the spurion field $\xi_{D_R}^\dagger$,

$$\langle \xi_{D_R}^\dagger \rangle = (0, 0, \tilde y_b)\, \Lambda\,, \tag{28}$$

in addition to (5). It breaks a sub-group of G_F',

$$SU(3)_{D_R} \times U(1)_T \to SU(2)_{D_R} \times U(1)_{\rm III}\,, \tag{29}$$

such that the unbroken flavour group is now given by

$$G_F'' = SU(2)_{Q_L} \times SU(2)_{U_R} \times SU(2)_{D_R} \times U(1)_{\rm III}\,. \tag{30}$$

(Alternatively, we could have assumed a large VEV for the spurion field $Y_D^{(2)}$. However, in this case one would have off-diagonal CKM elements of order 1 in contrast to observation.) The 10 unbroken generators of G_F'' are $T^{a=1,2,3}_{Q_L,U_R,D_R}$ and a charge operator for the 3$^{\rm rd}$ generation,

$$Q_{\rm III} = B - \frac{2}{\sqrt{3}}\, T^8_{Q_L=U_R=D_R}\,. \tag{31}$$

The additional 5 Goldstone bosons associated to the breaking of $G_F' \to G_F''$ now appear in the non-linear representation of the spurion field

$$\xi_{D_R}^\dagger = \langle \xi_{D_R}^\dagger \rangle\, \mathcal{U}^\dagger(\Pi_{D_R}) = \tilde y_b\, \Xi_{D_R}^\dagger\, \Lambda\,, \tag{32}$$

where $\mathcal{U}(\Pi_{D_R})$ and Ξ_{D_R} are defined analogously to (8,10). The reduction of the remaining spurion field $Y_D^{(2)}$ in terms of irreducible representations of G_F'' is given by

$$\begin{aligned} \tilde Y_D^{(2)} &\equiv Y_D^{(2)}\, \mathcal{U}_{D_R}^{(2)} \sim (2,1,2)_0\,, \\ \chi_L &\equiv Y_D^{(2)}\, \Xi_{D_R} \sim (2,1,1)_{-1}\,. \end{aligned} \tag{33}$$

In this way, the down-type Yukawa matrix is now parameterized as

$$\hat Y_D = \mathcal{U}(\Pi_{Q_L}) \begin{pmatrix} \tilde Y_D^{(2)}/\Lambda & \chi_L/\Lambda \\ 0\ \ 0 & \tilde y_b \end{pmatrix} \mathcal{U}^\dagger(\Pi_{D_R})\,, \tag{34}$$

which is to be diagonalized as

$$\hat{Y}_D = V_{d_L} \begin{pmatrix} y_d & 0 & 0 \\ 0 & y_s & 0 \\ 0 & 0 & y_b \end{pmatrix} V_{d_R}^\dagger \,, \qquad (35)$$

with two small eigenvalues $y_{d,s} = \mathcal{O}(\Lambda'/\Lambda)$. In particular, we have the relations

$$\begin{aligned} &\left(\mathcal{U}^\dagger(\Pi_{Q_L})\,\hat{Y}_D\,V_{D_R}\right)_{3j} = \tilde{y}_b \left(\Xi^\dagger_{D_R} V_{d_R}\right)_j \\ &= (\Xi^\dagger_{Q_L} V_{d_L})_{3j}\,(y_d)_j = V_{tb}\,y_b\,\delta_{3j} + \mathcal{O}(\Lambda'/\Lambda)\,, \end{aligned} \qquad (36)$$

and

$$|\tilde{y}_b|^2 = \sum_{j=1}^{3} V_{tj}\,(y_d^2)_j\,V_{tj}^* = y_b^2 + \mathcal{O}((\Lambda'/\Lambda)^2)\,. \qquad (37)$$

Considering again the example $b \to s\gamma$ in the case of large $\tan\beta$, the analogue of (20) is now given by a dimension-6 operator,

$$\mathcal{O}_{\rm eff} = \frac{1}{\Lambda^2}\,(\bar{Q}_L \Xi_{Q_L}\,\phi_D\,\sigma_{\mu\nu}\;\Xi^\dagger_{D_R} D_R)\,F^{\mu\nu} + \text{h.c.} \qquad (38)$$

After changing to the mass basis, using (19,36), we obtain

$$\left(V^\dagger_{d_L}\,\Xi_{Q_L}\,\Xi^\dagger_{D_R} V_{d_R}\right)_{ij} = V^*_{ti}\,V_{tj}\,(y_d)_j/\tilde{y}_b \simeq V^*_{ti}\,V_{tb}\,\delta_{3j}\,, \qquad (39)$$

which projects out the leading chirality structure in (21) for a large Yukawa coupling of the bottom quark [9].

The set of invariant quark bi-linears in the 2HDM with large $\tan\beta$ is constructed from building blocks which are singlets under $SU(2)_{U_R} \times SU(2)_{D_R}$, namely [10]

$$\begin{pmatrix} \bar{Q}_L\,\Xi_{Q_L} \\ \bar{U}_R\,\Xi_{U_R} \\ \bar{D}_R\,\Xi_{D_R} \end{pmatrix} \otimes \begin{pmatrix} \Xi^\dagger_{Q_L} Q_L \\ \Xi^\dagger_{U_R} U_R \\ \Xi^\dagger_{D_R} D_R \end{pmatrix}, \qquad (40)$$

and

$$\begin{pmatrix} \bar{Q}_L\,\mathcal{U}^{(2)}_{Q_L} \\ \bar{U}_R\,\mathcal{U}^{(2)}_{U_R} \\ \bar{D}_R\,\mathcal{U}^{(2)}_{D_R} \end{pmatrix} \otimes \mathcal{P}' \otimes \begin{pmatrix} \mathcal{U}^{(2)\dagger}_{Q_L} Q_L \\ \mathcal{U}^{(2)\dagger}_{U_R} U_R \\ \mathcal{U}^{(2)\dagger}_{D_R} D_R \end{pmatrix}, \qquad (41)$$

and

$$\begin{pmatrix} \bar{Q}_L\,\mathcal{U}^{(2)}_{Q_L} \\ \bar{U}_R\,\mathcal{U}^{(2)}_{U_R} \\ \bar{D}_R\,\mathcal{U}^{(2)}_{D_R} \end{pmatrix} \otimes \mathcal{P}' \cdot \chi_L \otimes \begin{pmatrix} \Xi^\dagger_{Q_L} Q_L \\ \Xi^\dagger_{U_R} U_R \\ \Xi^\dagger_{D_R} D_R \end{pmatrix}, \qquad (42)$$

with $\mathcal{P}'$ being an arbitrary polynomial of $Y_U^{(2)}\,Y_U^{(2)^\dagger}$, $\tilde{Y}_D^{(2)}\,\tilde{Y}_D^{(2)^\dagger}$, and $\chi_L\,\chi_L^\dagger$.

In summary, we have constructed an effective theory for flavour transitions beyond the SM, where the original (global) flavour symmetry G_F of the SM gauge sector is considered to be spontaneously broken by the large top – and, in the case of 2HDM with large $\tan\beta$, also bottom – Yukawa coupling. The associated Goldstone modes can be used to define a non-linear representation of flavour symmetry, where the concept of minimial flavour violation can be embedded by introducing spurion fields with canonical mass dimension. As a consequence, the dominance of certain flavour structures in rare quark decays – in particular, the top-quark dominance in flavour-changing neutral currents with down-type quarks – is a direct consequence of the $1/\Lambda$ expansion in the effective theory. The dynamical interpretation of the Goldstone modes as well as possible extensions beyond MFV (see, for instance, [7]) are left for future studies.

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[8] In the following, hatted quantities denote dimensionless Yukawa matrices, whereas the corresponding unhatted fields are spurions with canonical mass dimension.

[9] Notice that for $b \to s\gamma$ the VEV for ϕ_D in (38) provides an additional suppression factor $v_d/v = \cos\beta$.

[10] We do not discuss the issue of broken Peccei-Quinn symmetry [3] here. A detailed discussion can be found in [1].